

COSMIC RAYS FROM ACCRETING ISOLATED NEUTRON STARS

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ABSTRACT

Interstellar matter that is accreted onto isolated magnetic neutron stars in the Galaxy ($\sim 10^9$ by number) is accelerated and reflected back by MHD shocks, which envelope the stars. The integrated power in the Galaxy $L_{cr,ns}$ is $\gtrsim 10^{40} \text{ erg s}^{-1}$, the energy distribution is a power law of spectral index > 2 , and the particle energy can be raised to 10^6 GeV, consistent with the power and spectrum of primary cosmic rays in the Galaxy.

The major contribution for $L_{cr,ns}$ comes from a minority of $\sim 10^7$ isolated neutron stars which are located within dense clouds. Sources in these clouds, that are generally spread within the Galactic disk, can explain the concentration of high-energy cosmic rays in the Galactic plane, as deduced from pion decay spectra in gamma-ray observations. The soft X-ray luminosity from these neutron stars is consistent with the Galactic X-ray background. The accretion may be associated with ion-neutral bias, that is further enhanced by ion confinement in frozen-in magnetic fields, which can raise the relative abundance of first ionization potential (FIP) elements in the cosmic rays.

Key words: cosmic rays, acceleration of particles, shock waves, stars: neutron, X-rays: ISM

1 INTRODUCTION

The origin of cosmic rays has long been a central question in astrophysics (e.g. Cesarsky 1987, hereafter C87, Blandford & Eichler 1987 [BE87], Axford, 1994). The angular distribution of cosmic rays is highly isotropic. The isotropy accuracy is $\sim 0.1\%$ at energies $\lesssim 10^4$ GeV and $\sim 1\%$ at 10^8 GeV, but it is limited to a few tens of per cent in the ultra-high-energy range $> 10^9$ GeV. The origin of cosmic rays, apart of the ultra-high-energy class, is believed to be Galactic, associated with continuous leakage from the Galactic magnetic field confinement. The confinement time is observationally constrained by the decay time of unstable nuclei, e.g. ^{10}Be ($t_d = 2.2 \times 10^6$ yr), or ^{26}Al ($t_d = 0.85 \times 10^6$ yr), to be $\sim 8 \times 10^6$ yr (for homogeneous ISM).

Three major requirements of candidates for cosmic ray sources are that they produce the observed cosmic ray power, spectra, and elemental abundance. The total power of the Galactic component is estimated to be $L_{cr} \sim 2 - 3 \times 10^{40}$ erg s $^{-1}$ (BE87). The spectra, which are noisy and depend on the specific ion, are approximately a power law $\propto E^{-\alpha}$, where $\alpha < 2.5$ at energies below $\sim 10^2$ GeV, ~ 2.7 up to $10^{6.5}$ GeV and ~ 3 at $\gg 10^{6.5}$ GeV. The elemental composition (Meyer 1985a,b, Silberberg *et al.* 1987) at energies up to 10^6 GeV clearly deviates from the Local Galactic abundance, but fairly agrees with the Coronal composition. Specifically, there is a relative enhancement in cosmic rays of the low FIP (< 10 eV) elements.

Supernova remnants (SNRs) that heat and accelerate the surrounding material via shock wave acceleration are widely accepted as the origin of Galactic cosmic rays (Bell 1978, Blandford & Ostriker 1980, BE87). The cosmic ray power is only a few per cent of the total power associated with SNR mass motions, therefore, even an inefficient shock acceleration mechanism is applicable. Such a mechanism also explains the spectral shape. The maximal energy gain per particle E_{max} , however, is limited by the shock lifetime, estimated to be $\sim 10^5 B / (10^{-6} \text{G}) \text{GeV}$ in optimal conditions (E_{max} may be increased to 5×10^6 GeV if the SNR is propagating into a stellar wind cavity [Völk & Biermann 1988]). Moreover, the FIP bias is unexpected, since the SNR is accelerating material which is believed to have the Local Galactic abundance (this problem is relaxed in models of SNR in a stellar wind cavity, if the presupernova object is a red or a blue giant, or a Wolf-Rayet star [Silberberg *et al.* 1990]). In competing 'stellar' models sub-relativistic particles are injected from F and G stars (Meyer 1985a,b), which explain the cosmic ray composition, but require a secondary stage of acceleration, e.g. by passing a supernova shock wave. An injection mechanism for ultra-high-energy cosmic rays was proposed by Kazanas & Ellison (1986), based on accreting binary systems such as massive X-ray binaries, where a strong stellar wind of the companion star is accreted onto the compact object, a neutron star (NS) or a black hole.

Cowsik & Lee (1982) have shown that accretion flows with shocks onto isolated NSs can accelerate cosmic rays very efficiently. They consider old NSs, assuming that their surface magnetic fields are weak and do not affect the hydrodynamics. The flow stagnation imposed by the stellar surface results in a standing spherical shock wave which envelops the star very close to its surface. Potential energy, associated with interstellar matter

(ISM) accreted onto the star, is continuously channeled to acceleration of high energy particles. The cosmic ray power is $\sim 10^{36}$ erg s $^{-1}$ per NS, assuming that its velocity is smaller than the ISM sound speed, and the ISM density is large $\sim 10^3$ cm $^{-3}$. For such a luminosity a few times 10^4 NSs are sufficient to explain the Galactic cosmic rays. Such a high power, $\sim 10^{-2}$ of the Eddington luminosity, may be, however, associated with a temperature increase above the pair production threshold, as was recently shown by Turolla *et al.* (1994). Pair annihilation and gamma ray emission can significantly cool the particles and (electron) synchrotron would become very efficient, whether the magnetic fields are originated in the NS or frozen within the accreted ISM. Also, if the peculiar velocity of an old NS is larger than the ISM sound speed the available power is significantly smaller than 10^{36} erg s $^{-1}$. More importantly, various observations suggest that the magnetic field of an old NS remains very strong (e.g. Srinivasan 1989, Phinney & Kulkarni, 1994), and in that case the magnetic pressure can halt the flow at the Alfvén surface far from the stellar surface.

In this paper I focus on Galactic cosmic rays that originate from accreting magnetic isolated NSs. Because of the halting of the accretion, a standing shock is created, within which shock acceleration occurs (Shemi 1995). The available power, compared with the free fall luminosity, is significantly reduced, to a value orders of magnitude below the Eddington luminosity. Other mechanisms that might regulate the accretion, such as the "propeller effect" (Illarionov & Sunyaev 1975) or a relativistic wind induced by the NS, which will not be discussed here. Accretion through the polar caps, which, unlike disk-accretion, probably is inefficient in spherical inflows, is ignored, and instabilities of the magnetosphere are not considered. The accreted matter is never heated to ultrarelativistic temperature, and pair production therefore does not take place.

The total population of Galactic isolated NSs is estimated to be $N \sim 10^9$ (Blaes & Madau 1993, hereafter BM93), but the major contribution to cosmic rays comes from a small subgroup ($\sim 1\%$) containing the slow NSs which are located within dense clouds. In the following paper I first describe the physics of accretion, shock waves and particle acceleration around an isolated magnetic NS, and then discuss the applicability of these objects as sources of cosmic rays.

2 ACCRETION ONTO A MAGNETIZED NEUTRON STAR

An isolated NS that is propagating with velocity $v_{40} = v/(40 \text{ km s}^{-1})$ accretes ISM from a cylinder of a Bondi-Hoyle radius

$$r_{BH} = \frac{2GM}{v^2 + c_s^2} \simeq 2.3 \times 10^{13} v_{40}^{-2} \text{ cm} . \quad (1)$$

Here c_s is the sound speed and M is the NS mass, where throughout this paper we use $M = 1.4M_\odot$. I will focus on dense ISM clouds ($n = 10^4 n_4 \text{ cm}^{-3}$), where the accretion rate would be large. Extreme UV and soft X-ray radiation emitted by the accreted plasma is sufficiently intense to ionize the matter within a Strömgren sphere $r_S \sim 3 \times 10^{15} (S/10^{43}) n_4^{-2/3} (\alpha_B/10^{-12})^{-1/3} \text{ cm}$, where S is the rate of ionizing photons,

in photons s^{-1} , and α_B is the recombination coefficient to excited states of hydrogen, in $\text{cm}^3 \text{s}^{-1}$. In the ionized gas macroscopic magnetic fields and collective plasma phenomena couple the particles, and the accretion become hydrodynamical, rather than collisionless. Removal of angular momentum becomes efficient and raises the accretion efficiency, by a factor as large as $(c/c_s)^2 \sim 10^9$, compared with that of a neutral gas (e.g. Shapiro & Teukolsky 1983).

Far from the central object ($r_* \ll r \lesssim r_{BH}$, where r_* is the NS radius, 10^6 cm in this paper) irreversible processes of energy dissipation are negligible, and the pressure obeys the adiabatic relation

$$p = p_0 \left(\frac{\rho}{\rho_0} \right)^\Gamma, \quad (2)$$

where Γ is the specific heat ratio. The expected accretion rate is

$$\dot{M} = 4\pi r_{BH}^2 \lambda_s \rho v \simeq 4.3 \times 10^{14} n_4 v_{40}^{-3} \text{ g s}^{-1}, \quad (3)$$

where λ_s is a numerical factor of order of unity. The accretion rate is much smaller than the Eddington rate $\dot{M}_{Ed} = L_{Ed} r_*/(GM) \sim 10^{18} \text{ g s}^{-1}$, therefore radiation pressure is never strong enough to halt the accretion.

The continuous removal of angular momentum causes streamlines to bend and converge toward the NS. To follow the flow dynamics we neglect the influence of magnetic fields and back streams from the central object or from the shock, and assume spherical symmetry. The dynamics are governed by the continuity equation (3) and by the Bernoulli equation

$$\frac{v^2}{2} + \frac{c_s^2}{\Gamma - 1} - \frac{GM}{r} = 0. \quad (4)$$

The sound speed [$c_s = (\gamma kT/m)^{-1/2} = 0.57 (T/50 \text{ K})^{1/2} \text{ km s}^{-1}$] of the dense and cold ISM cloud is low, therefore, even for small NS velocity, the flow at $r = r_{BH}$ is supersonic, with a rather large Mach number $\bar{M} = 70 v_{40} T_{50}^{-1/2}$. Unless nonthermal acceleration occurs, the energy gain per particle does not exceed $[1 - (v_{ff}/c)^2]^{-1/2} - 1 = 0.34$ of its rest mass, maintaining the plasma subrelativistic even when approaching the NS, with typical blackbody temperature $\sim 10^6 - 10^7 \text{ K}$. Substituting the nonrelativistic value $\Gamma = 5/3$ and the limit $c_s \ll v$, equations 2-4 give the scaling behavior of v , c_s , ρ , and the gas ram pressure ρv^2 :

$$v \propto r^{-1/2}; \quad c_s \propto r^{-1/2}; \quad \rho \propto r^{-3/2}; \quad \rho v^2 \propto r^{-5/2}. \quad (5)$$

Note that the Mach number remains constant. The ionized gas has high conductivity and the magnetic Reynolds number is large, so that the residual magnetic fields in the ISM are frozen within the inflow. Field lines are swept along the streamlines and the global field pattern becomes radial [one may consider a large - scale field which initially is homogeneous $(B_x, B_y, B_z) = (B_0, 0, 0)$, in where an accreting object is placed. The lines of the frozen-in field would be bent towards the central object, along the y and z axis, resulting in a configuration of a nozzle. The nozzle configuration is narrow in its central part and globally the field lines surrounding the central object become radial]. The field strength increases $\propto r^{-2}$ and the magnetic pressure would therefore increase as r^{-4} , faster than the ram pressure and the gas pressure ($\Gamma^{-1} \rho c_s^2 \propto r^{-5/2}$). Such a field enhancement

would, however, be regulated when small scale magnetic reconnections become sufficiently fast. The field strength is therefore bounded by the equipartition limit, that is achieved when the reconnection time scale, r/v_A , becomes comparable to the compression time scale, r/v_{ff} , $B_{gas}^2/(4\pi)^{-1} \lesssim \rho v^2$. This occurs at the equipartition radius

$$r_{eq} = 7 \times 10^8 \left(\frac{B_{ISM}}{3 \times 10^6 \text{ G}} \right)^{4/3} n_4^{-2/3} v_{40}^{-10/3} \text{ cm} \quad (6)$$

where B_{ISM} is the magnetic field strength at r_{BH} .

Magnetic fields of the strength expected in old NS, $10^9 \lesssim B_{\text{surface}} \lesssim 10^{12} \text{ G}$ (e.g. Srinivasan 1989, Phinney & Kulkarni 1994), would cause stagnation of the flow at the Alfvén surface, where the magnetic pressure balances the flow ram pressure. The flow is decelerated and compressed, and standing shock waves occur ahead of the Alfvén surface.

Using the dipole approximation $B \propto r^{-3}$, the Alfvén radius is given by

$$r_m = 10^8 B_{10}^{4/7} n_4^{-2/7} v_{40}^{6/7} \text{ cm} . \quad (7)$$

Here B_{10} is the NS surface magnetic field in units of 10^{10} G .

If the light corotation radius $r_{lc} = cP/2\pi = 4.8 \times 10^9 P_1 \text{ s}$ ($P_1 \equiv P/1 \text{ s}$) is larger then r_m a pulsar wind beyond the light cylinder can halt the inflow. However, such a mechanism would be significant if the NS rotation period is shorter than

$$P_{lc,m} = 2 \times 10^{-2} B_{10}^{4/7} n_4^{-2/7} v_{40}^{6/7} \text{ s} , \quad (8)$$

while old NSs, unless recycled, are generally slower than $\sim 1 \text{ s}$. Beyond r_{lc} the magnetic field decreases as r^{-1} , therefore in very fast rotators ($P < P_{lc,m}$) the Alfvén surface would generally be pushed beyond r_{BH} . If $r_m > r_{BH}$, the accretion is prevented in the first place. The motion of the NS through the ISM is then associated with a stationary bow shock, distant

$$r_{ism} = 3.2 \times 10^{13} B_{10} v_{40}^{-1} P_1^{-2} \text{ cm} \quad (9)$$

ahead of the star. The accretion flow will overcome the centrifugal barrier (where the free fall velocity become equal to the Keplerian velocity), at the centrifugal radius $r_{cent} = (GM/\Omega^2)^{1/3} = 1.7 \times 10^8 P_1^{2/3} \text{ cm}$, if the star period P exceeds $2\pi r_m^{3/2} (GM)^{-1/2}$. With equation 7 this condition reads

$$P > P_{cent,m} = 0.46 B_{10}^{6/7} n_4^{-3/7} v_{40}^{9/7} \text{ s} . \quad (10)$$

For NSs in dense ISM clouds we can reasonably assume $P > P_{cent,m}$, and therefore consider a shock wave at $r_m \sim 10^8/r_{m,8} \text{ cm}$.

3 PARTICLE ACCELERATION

The topology and physics of the shock ahead of the Alfvén surface is complex and not fully

resolved yet. Shock topology would differ from a spherical shape, presumably is conical with an opening angle $\sim \bar{M}^{-1}$ and a bow front, and the obliquity varies from parallel to perpendicular. A large electron and ion concentration is built up in the foreshock regions. Irregularities in the shock topology and temporal variability are also expected, in the same manner as in simulations of wind accretion onto an *unmagnetized* massive object (Ruffert, 1994 and references therein). Backward streams and perturbations influence the topology, sometimes depending on the Mach number and the r_{BH}/r_* ratio, making the flow configuration unstable.

The jump conditions for a planar-discontinuity, infinitesimally narrow shock imply a compression ratio $\zeta = \rho_2/\rho_1$ (here '1' denotes the upstream region and '2' stands for downstream) that approaches the limiting value $\zeta \rightarrow (\Gamma + 1)/(\Gamma - 1) = 4$ for $\Gamma = 5/3$ (e.g. Draine & McKee 1993). The parallel component of the magnetic field is conserved ($B_{1,\parallel} = B_{2,\parallel}$), while the perpendicular component downstream is enhanced ($B_{2,\perp} = \zeta B_{1,\perp}$ for a flow with no transverse velocity component on both sides). The true strength of the magnetic fields upstream is reduced if the NS dipole fields are screened by electric field ($\bar{v} \times \bar{B}$) in the shock region. On the other hand, the enhancement of frozen magnetic fields in the inflow would increase the upstream field.

In the following I will neglect the complexity in the shock conditions, and assume that known acceleration mechanisms in planar shocks are still applicable, to assess the cosmic ray luminosity, the maximal energy per particle and the energy distribution.

Some energy loss mechanisms take place at the shock region, among them radiation, dissipation through MHD waves and escape of neutral particles. Defining f_{cr} the fraction of the free fall luminosity ($l_{ff} = \dot{M}GM_m^{-1}$) that is left in cosmic rays after such losses, we obtain the power per an accreting NS

$$l_{cr,ns} = f_{cr} \dot{M} \frac{GM}{r_m} = 8 \times 10^{32} f_{cr} \lambda_s B_{10}^{-4/7} n_4^{9/7} v_{40}^{-15/7} \text{ erg s}^{-1} . \quad (11)$$

Particles would be accelerated through multiple shock crossing, presumably via a *diffusive shock* first-order Fermi mechanism (Drury 1983, BE87). *Shock drift* acceleration, moreover, can take place in the quasi perpendicular segments of the shock, where an electric field $\bar{v} \times \bar{B}$ exists. The nuclear mean free path $l_{col} = \lambda_{pp}/nm_p$ (where $\lambda_{pp} \sim 60 \text{ g cm}^{-2}$) at the shock region is $\sim 3 \times 10^{13} n_4^{-1} \text{ cm}$, much larger than the shock scale r_m . However, the proton Larmor radius $l_B = \gamma m_p c^2 / (eB)$ will be smaller than r_m unless particles are ultra relativistic ($\gamma \gg 10^5$). Particle - wave interactions will therefore dominate the inelastic (particle - particle) collisions, making the shock collisionless. Particles are scattered by Alfvén waves, small-scale magnetic irregularities and turbulences, or other collective phenomena.

If the upstream magnetic pressure P_{mag} is well below the gas pressure P_{gas} , the MHD perturbations are basically Alfvén waves that travel upstream with velocity $v_A < v_1$. In the downstream region the plasma experiences strong magnetic turbulences. Particles are reflected by these turbulences backward to the upstream region, then reflected again downstream by the Alfvén waves (or also by turbulences, if $P_{mag} \sim P_{gas}$ upstream. Note that even in that case we expect a large pressure gradient which would force a shock). For a particle with velocity w the momentum gain per single shock crossing is $p \rightarrow p \times$

$4[v_2 - v_1]/3w$. After a sequence of $\sim c/w$ random crossings the momentum is raised to the relativistic regime, and, since $w \gg v_A$, the Alfvén waves upstream are seen as a quasi-static barrier.

When the acceleration is sufficiently rapid, and not affected by energy losses, the particle energy is raised until the particles diffuse from the shock. The conical shape of the shock allows downstream accelerated particles to escape. The diffusion coefficient may be assumed to have the general form $\kappa = \lambda w/3$, where λ is the particle mean free path, usually between one to a few Larmor radii. Particles remain confined within the shock as long as the upstream diffusion length scale κ/v_1 is larger than the shock radius of curvature $\sim r_m$. The maximal energy per particle can be estimated by using the Bohm approximation ($\lambda = l_B$) and requiring $(l_B w)/(3v_1) < r_m$, that gives (with $w = c$; $v_1 = (2GM)^{1/2} r_m^{-1/2}$)

$$E_{max} \equiv \gamma_{max} m c^2 \sim 5.4 \times 10^4 r_{m,8}^{-5/2} B_{10} \text{ GeV} . \quad (12)$$

To obtain a value of E_{max} comparable to the 'knee' value $\lesssim 10^6$ GeV, at which the cosmic ray spectrum bends, one needs $r_{m,8} \gtrsim 0.3$. Since the Alfvén radius scales as $r_m \propto B^{4/7}$, E_{max} scales $\propto B^{-3/7}$, and such large values of E_{max} does not have to be associated with a large accretion rate, if the stellar magnetic field is small $B_{10} \ll 1$.

Energy loss by synchrotron emission also can, in principle, regulate the particle acceleration. The maximal energy in that case is determined by the competition between these processes. The time scale for diffusive acceleration $t_{acc} = p/\dot{p} \sim \kappa/v_1^2$ [more rigorously, taking care for the time spent by a particle in each side of the shock, one can obtain $t_{acc} = 3 \times (v_1 - v_2)^{-1} (\kappa_1/v_1 + \kappa_2/v_2)$, e.g. Drury 1983] while the ion cooling time scale $t_{synch,ion} \sim 6 \times 10^{18} A^3 Z^{-4} B^{-2} \gamma^{-1}$ s (here B in Gauss). From the requirement $t_{synch,ion} > t_{acc}$ we obtain the constraint $\gamma_{max,ion} \lesssim 2.6 \times 10^8 A^{3/2} Z^{-2} B_{10}^{-1/2} r_{m,8}$ (note that in this case $\gamma_{max,ion}$ only weakly depends on B , namely $\gamma_{max} \propto B^{1/14}$).

The synchrotron cooling time of electrons is, however, shorter than that of protons, by a factor $(m_e/m_p)^3$, implying (with $\gamma_{max,e} \propto (m_e/m_p)^{3/2}$) $\gamma_{max,e} \sim 3.3 \times 10^3 B_{10}^{-1/2} r_{m,8}$ s. Synchrotron cooling although negligible for ions below $\sim 10^8$ GeV, becomes acute for electrons. Moreover, electrons acceleration is not even started unless additional 'injection' mechanism operates. MHD waves upstream do not resonate with electrons below $\gamma_{min,e} - 1 = m_p/m_e \times (v_A/c)^2$ (for $(c/v_A)^2 \gg (m_p/m_e)^2$, hence only the relativistic electrons can be confined in the shock. The mechanism discussed here therefore may not be applicable for cosmic ray electrons.

The initial energy distribution of particles upstream is roughly a delta function at $GM/r_m \sim (0.064c)^2 r_{m,8}^{-1}$, with some thermal broadening ($kT \sim c_s^2 \ll v_1^2$). By the stochastic nature of the acceleration, and since the total number of scatterings per particle is sufficiently large, the distribution is spread out with no preferred scale. From the linear theory of acceleration in subrelativistic shocks we expect a general form

$$n(E) \propto E^{-\alpha} . \quad (13)$$

The spectral index, given by

$$\alpha = \frac{\zeta + 2}{\zeta - 1} \sim 2 , \quad (14)$$

is independent of the magnetic fields, and is found to be typical in other cases of shocks. Cowsik & Lee (1982) show that a power law typical to cosmic rays can also be obtained in shock waves which have a spherical geometry. Ellison & Eichler (1985) also show that $\alpha \gtrsim 2$ in non-linear shock models that account for the cosmic ray feedback. Even when the Alfvén Mach number gives rise to strong compression, $\zeta \gg 4$, the slope is almost unaffected.

4 DISCUSSION

The Galactic isolated NS population (BM93) is estimated from the pulsar birthrate, assuming a steady state, to be $N > 10^8$, but nucleosynthesis constraints on Galactic chemical evolution require a total number as large as $N \sim 10^9$ (Arnett, Schramm, & Truran 1989). A fraction of the potential energy associated with matter accreted onto these NSs is continuously channeled to accelerate particles to very high energies. Although the detailed shock structure and acceleration process are yet unclear, the energy budget and the knowledge of acceleration in astrophysical MHD shock waves suggest that accreting isolated NSs are significant cosmic ray sources. The cosmic ray luminosity $L_{cr,ns}$ is the integral over the cosmic ray emission $l_{cr,ns}$ from all the Galactic isolated NSs, and other Galactic compact accretors like isolated white dwarfs, black holes, and massive binaries, that should be taken into account.

We briefly discuss the parameters that affect the accretion rate and thereby the cosmic ray luminosity, namely the ISM density distribution, the NS velocity distribution, their surface magnetic fields and periods, and the physics behind the efficiency factor f_{cr} .

The filling factor of the ISM dense phase is estimated to be $\sim 1\%$ in the disk, and it rises to $\sim 10\%$ in an extended region in the Galactic center. From the estimated NS density in clouds $5 \times 10^{-4}(n_{ns}/\text{pc}^{-3})(M_{cloud}/M_{\odot})^{3/2}$, BM93 deduced $\sim 10^3$ isolated NSs in the 19 clouds closest (≤ 800 pc) to the Solar System.

The rms value of their velocity distribution is calculated to be 98 km s^{-1} , where a substantial fraction of them, $f_s \sim 0.25$, are slow ($v \lesssim 40 \text{ km s}^{-1}$), and $f_d \lesssim 0.2$ of them are located within dense ISM clouds ($n \sim 10^2 - 10^8 \text{ cm}^{-3}$). The upper density value represents the cores of giant molecular clouds. BM93 show a disk velocity distribution with 6% of NSs having $v \leq 20 \text{ km s}^{-1}$, 22% below 40 km s^{-1} , and 50% below 72 km s^{-1} . It turns out that the major contribution to $L_{cr,ns}$ would come from a minority of 10^7 isolated NSs which are slow and located within dense clouds.

The strengths of magnetic fields of old NS are controversial. Field decay is predicted in many NS models, which have also attempted to explain why such decay slows down or stops. However, there are arguments against a decay mechanism, in which case the diversity in pulsar magnetic fields is the result of the conditions during their creation (see e.g. Srinivasan, 1989 Phinney & Kulkarni 1994).

The fraction f_{cr} of the input power that is channeled to cosmic rays depends on the efficiency of competing energy loss mechanisms. Synchrotron radiation can maintain the electron temperature low but only loosely affect the ions. Free - free emission [$\epsilon_{ff} \times r_m^3 \sim 10^{27} n_4^2 (T/10^7 \text{ K})^{1/2} \text{ erg s}^{-1}$] and Compton losses are also small. Escape of neutrons,

which are produced through $pp \rightarrow nX$ reactions or photodissociation of ^4He nuclei, can significantly ($\lesssim 25\%$) affect the ultra-high-energy part of the spectrum (Kazanas & Ellison 1986). Energy is also dissipated through MHD waves. For the sake of evaluating $L_{cr,ns}$ we will use $f_{cr} \sim 0.25$ and focus on the slow stars in dense clouds. Using the canonical values discussed above we have

$$L_{cr,ns} = 1.2 \times 10^{40} \left(\frac{N}{10^9}\right) \left(\frac{f_s}{0.25}\right) \left(\frac{f_d}{0.2}\right) \left(\frac{f_{cr}}{0.25}\right) \left(\frac{L_{ff}}{10^{33}}\right) \text{ erg s}^{-1}, \quad (15)$$

which is consistent with the power estimated for Galactic cosmic rays.

The maximal energy gain per particle (Equation 12) is sufficient to explain the majority of cosmic ray particles, but not the very-high-energy ones. With a large accretion rate E_{max} would greatly be raised, to values similar to those proposed by Kazanas & Ellison (1986) for binary systems. Moreover, coherent acceleration through magnetic field annihilation (frozen-in fields + NS fields) as is suggested here, as well as pulsar acceleration, as proposed elsewhere by many authors (e.g. Srinivasan 1989, Arons & Tavani 1994), are likely to take place.

Of particular interest is the subgroup of very slow ($v \lesssim c_s$) isolated NSs in dense and cold clouds $c_s < 1 \text{ km s}^{-1}$. Although these strong accretors emit a large X-ray luminosity, their Bondi - Hoyle radius $\gg 10^{16} \text{ cm}$ can exceed the Strömgren radius. The *effective* accretion radius, where the gas is coupled through magnetic fields and angular momentum is removed by fast collisions, shrinks below r_{BH} . Accreting NSs with $r_s < r_{BH}$ may enlarge the cosmic ray (low FIP)/(high FIP) elemental ratio, that is observed to be ~ 3 times larger than in Local Galactic matter (Meyer 1985). Selection between ions and neutral atoms occurs far from the NS, at the radius where each element becomes ionized. The radii of photoionization fronts would increase as the ionization potential decreases (note that the ionized continuum is much flatter than blackbody). By assumption, ions are more effectively accreted than neutrals, and by such an ion - neutral bias the relative fraction of low FIP elements in the inflow is enhanced. Furthermore the ions would be anchored within the inflow by the frozen magnetic fields, while neutral atoms can diffuse in a direction oppose to the temperature gradient, As long as the diffusion time scale r^2/D is smaller then the dynamic time scale, r/v , ion-neutral separation by diffusion can be significant. The diffusion coefficient for neutrals in an ionized hydrogen gas (Geiss 1982) $D = 1.1 \times 10^{21} T / (10^4 \text{ K})^{3/2} n_1^{-1} \text{ cm}^2 \text{ s}^{-1}$ may, indeed, be larger than $r_{BH} v \sim 10^{18} - 10^{19} \text{ cm}^2 \text{ s}^{-1}$. The ion-neutral bias is particularly applicable in dense clouds $n_4 \gtrsim 1$ or, more precisely, when $n_4^{1/3} v_{40}^{-1} \gtrsim 70$, otherwise the accretion radius is well within the ionization sphere and the inflow is ionized in the first place.

Gamma-ray observations show a sharp concentration of high energy radiation in a very thin Galactic disk, with evidence for imprints of the Galactic arms. If this emission comes from decaying pions, originating in cosmic ray interactions with the ambient gas (Stecker & Stecker, 1970), it requires the cosmic ray concentration to follow the high energy photons. In our model this is fulfilled since the dense ISM clouds are also concentrated in the spiral arms, generally close to the galactic equator.

What is the electromagnetic imprint of isolated NSs? The soft X-ray luminosity from the shock region is a fraction of $\propto L_{cr}(1 - f_{cr})$, and the expected spectrum is nonthermal.

From equations 13 and 14 one can predict a power law of spectral index $(\alpha - 1)/2 \sim 1/2$. Other emission mechanisms, involved with heat loss of the star itself would basically produce thermal emission, but a magnetized atmosphere could affect the radiation transfer and modify the blackbody curve. A crucial point to note is that the Galactic diffuse X-ray background hardly restricts the cosmic - ray luminosity from isolated NSs. The integrated flux from isolated NSs was found to contribute no more than a fraction of a per cent to the soft X-ray background (BM93). The strongest accretors are located in dense clouds, and so even if they emit strong soft X-ray radiation, it would be greatly absorbed. Using a photoionization cross section $\sim 10^{-22}(E/\text{KeV})^{-3} \text{ cm}^2 (\text{H} - \text{atom})^{-1}$, the cloud optical depth $\tau_x \sim 10^2(E/\text{KeV})^{-3} n_4(R/10^{20}\text{cm})$, where R is the cloud radius. Solitary isolated NSs, moreover, may be the origin of the enhanced X-ray emission within $\pm 40^\circ$ longitude of the Galactic center, as observed by *EXOSAT* (BM93, Maoz & Grindlay 1994).

Although isolated NSs are the majority of the NS population, they have so far received only limited attention compared with that given to the brighter ones, the pulsars and the (NS) X-Ray sources. Only a small fraction of the Galactic isolated NSs can be observed directly. Following the early suggestion that nearby isolated NSs might be detectable due to radiation released from the accretion of the ISM (Ostriker, Rees & Silk 1970), Madau & Blase (1994) argue that the ~ 100 sources detected in the EUV - soft X-ray range the *ROSAT* - *WPC* and *EUVE* all - sky surveys, and which have no previous identification in the optical or other energies, are the natural candidates for isolated NSs. Recent observations at the Wise Observatory of the fields of ~ 15 of these sources show that indeed in many cases the sources are likely hot or peculiar, and unlikely to be white dwarfs or late type flare stars. Simultaneous observations planned from a number of experiments aboard the *Spectrum X Gamma* observatory, in UV (*TAUVEX*), EUV (*EUVITA*), and X-Ray bands (*SODART*, *Jet-X*, *MART*), will be extremely useful for the identification of such objects. The predicted characteristics to be studied in such UV - X-ray observations include spectral index, the presence of synchrotron absorption lines, temporal modulation, proper motion, and HII cometary - like regions. Detection of isolated NSs would be an important step towards our understanding of this Galactic component. The model presented here can be constrained by observations of the total isolated NS population, their spatial and velocity distributions, and the fraction of them in dense clouds.

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